

Neutrino Mass and Proton Decay in a $U(1)_R$ Symmetric Model

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Abstract

We study a $U(1)_R$ symmetric extension of supersymmetric standard model with supersymmetry breaking in the visible as well as hidden sectors. Specifically we study $U(1)_R$ breaking effects parametrized by the gravitino mass. A special R -charge assignment of right-handed neutrinos allows us to have neutrino Yukawa couplings with the R -charged Higgs field, which develops a tiny vacuum expectation value after the inclusion of $U(1)_R$ symmetry breaking. Even with $O(1)$ Yukawa couplings, a suitable size of Dirac neutrino masses can be generated if the gravitino mass is very small, $m_{3/2} = 1\text{--}10\text{ eV}$. Our flipped R -charge assignment also allows a new type of dimension five operator that can induce the proton decay. It turns out that the proton stability mildly constrains the allowed range of the gravitino mass: Gravitino heavier than 10 keV can evade the proton decay constraint as well as cosmological ones. In this case, the largest neutrino Yukawa coupling is comparable to the electron Yukawa. We also calculate the mass of the pseudo goldsino and its mixing to neutralinos, and briefly discuss its implications in cosmology and Higgs phenomenology.

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I. INTRODUCTION

Weak scale supersymmetry (SUSY) has been an attractive candidate for the physics behind the electroweak symmetry breaking (EWSB). SUSY breaking of order of the weak scale can trigger the desired electroweak symmetry breaking either at tree level or via radiative corrections. However, null results at LHC for SUSY particle search up to now would require us to reconsider such picture, providing a motivation of extending the minimal SUSY Standard Model (MSSM).

$U(1)_R$ -symmetric extension of the SUSY standard model is an interesting starting point for the physics beyond the MSSM. A nice feature of $U(1)_R$ symmetry in the matter sector is that it naturally explains the absence of baryon number violating operators that would lead to fast proton decay. See Ref. [1] for a solution to SUSY flavor problem. The gauge sector can be made $U(1)_R$ symmetric if each gaugino has Dirac mass, instead of the usual Majorana one. Such Dirac mass term can be generated from hidden sector SUSY breaking through a supersoft operator [2], which induces finite soft scalar masses. Models with Dirac gaugino have an advantage of reducing the degree of fine tuning in the Higgs potential even when the colored sparticles are as heavy as multi TeV [3]. The Higgs sector can also be made $U(1)_R$ symmetric if we introduce the mirror partner, R -partner of the MSSM Higgs doublets.

The origin of the EWSB can be addressed in a $U(1)_R$ -symmetric manner. Ref. [4] proposed¹ a natural realization of the EWSB by coupling the $U(1)_R$ -symmetric Higgs sector to visible sector SUSY breaking, in which the supertrace sum rule is avoided by the presence of the hidden sector SUSY breaking. Such coupling can be used to raise the lightest Higgs mass. The latter point is important if Dirac gaugino mass terms are generated by the supersoft operators which also suppress the tree level D -terms, as was noted in Ref. [2]. See also Ref. [5] and references therein. Implications in cosmology as well as Higgs search were also discussed in Ref. [6], in which it was pointed out that the visible SUSY breaking with Majorana gauginos is cosmologically disfavored.

In the present paper, we are interested in yet another notable aspect of R -symmetric extension of SUSY standard model. The $U(1)_R$ symmetry should be broken in the hidden sector for the cosmological constant to be canceled. If the R -symmetry breaking is mediated to the visible sector in a minimal way, an R -partner Higgs field develops a tiny vacuum expectation value (VEV) characterized by the gravitino mass. We will show that such a tiny VEV can be related to the smallness of neutrino masses. The generation of neutrino mass from SUSY breaking effects was studied before in Refs. [7, 8]. The generation of (Majorana) neutrino mass from R -symmetry breaking was also discussed recently in Ref. [9, 10]. In the present paper, we will relate the tiny R -breaking VEV to Dirac neutrino masses by flipping the R -charge assignment of the right-handed neutrinos. Our flipped assignment of $U(1)_R$ charges allows us to write a desired coupling of the right-handed neutrino to the R -partner Higgs doublet; at the same time, it also allows a new type of dimension five operator that

¹ One of the motivations behind the construction in Ref. [4] was to explore the possibility of testing the SUSY breaking mechanism through the Higgs sector as a portal. Here we are interested in a simple realization of EWSB via the visible SUSY breaking.

makes the proton unstable. It turns out that the constraint from the proton stability is milder than the corresponding constraints in the MSSM. Thus in our flipped $U(1)_R$ model, tiny neutrino masses can be generated without spoiling so much a nice feature of $U(1)_R$ symmetry of ensuring proton stability.

The present setup of $U(1)_R$ -symmetric model has potentially interesting implications on Higgs phenomenology. Since we consider the model of visible SUSY breaking and also assume that Dirac gaugino masses are induced from another SUSY breaking in a hidden sector, there appears a physical pseudo goldstino state ζ [11], which directly couples to the Higgs sector. As discussed in Ref. [6], such pseudo goldstino can affect the Higgs decay modes. We perform a similar analysis in our $U(1)_R$ symmetric setup and examine to some details the mass and the mixing of the pseudo goldstino. It turns out that although the pseudo goldstino can get a mass due to R -breaking effect, its mass m_ζ is quite suppressed compared to the gravitino mass $m_{3/2}$. This is contrasted to the ‘sequestered’ case in which m_ζ is twice as large as $m_{3/2}$ [11, 12]. Moreover, the mixings of the pseudo goldstino to the MSSM Higgsinos and gauginos are highly suppressed due to the softly broken $U(1)_R$ symmetry.

Another important constraint on the present model comes from cosmology; the light gravitino is constrained rather severely not to disturb the successful Big-Bang Nucleosynthesis (BBN), the structure formation of galaxies and the cosmic microwave background (CMB) radiations, and not to be produced too much in the early universe. As we shall see later, the allowed range of the gravitino mass is given by²

$$m_{3/2} \lesssim 16 \text{ eV} , \quad 10 \text{ keV} \lesssim m_{3/2} . \quad (1)$$

We will take account of these limits in the following analysis. We will also make a brief comment on the limit on decaying dark matter scenario from the diffused gamma-ray line search.

The paper is organized as follows. In the next section, we present our model first in the R -symmetric limit, and we then include R -breaking effects in a minimal way that are parameterized by the gravitino mass. Throughout the present paper, we assume the gravitino mass to be much smaller than the weak scale, although we do not elucidate possible origins of hidden sector SUSY breaking. In section three, we show how the small VEVs of the R -charged Higgses can be used to explain the smallness of the neutrino mass. We also discuss the constraint from the proton decay in our flipped $U(1)_R$ model for neutrino mass. In section four, we calculate the pseudo goldstino mass and the mixings to neutralinos under the assumption of minimal $U(1)_R$ breaking. Specifically we show how the mass of the pseudo goldstino is suppressed compared with the gravitino in the combined model of Dirac gaugino and visible SUSY breaking. We then apply the results to briefly discuss implications of the present model on Higgs phenomenology and cosmology. The final section is devoted to our conclusion and discussion. We also add two appendices concerning the analysis on baryon and/or lepton number violating operators and a possible extension of the model in which the constraint from proton stability will be relaxed via $U(1)$ flavor symmetry.

² See Ref. [13] and references therein. See also Sec. IV C for a brief discussion.

For simplicity, we will often refer to $U(1)_R$ symmetry as “ R symmetry”, if no confusion is expected.

II. R -SYMMETRIC MODEL WITH VISIBLE SUSY BREAKING

In this section, we present an extension of the MSSM, in which $U(1)_R$ symmetry is realized by combining the model of Dirac gauginos [2] with the model of visible SUSY breaking of Ref. [4]. We then discuss the effects of the minimal R -symmetry breaking induced from the coupling to supergravity.

A. $U(1)_R$ Symmetric Model

The $U(1)_R$ symmetry forbids the usual Majorana gaugino mass terms as well as the higgsino mass term, $\mu H_u H_d$, in the MSSM. To realize $U(1)_R$ symmetry, we extend the MSSM by introducing the R -partners of the gauginos and those of the higgsinos.

First let us briefly discuss the gauge sector. Let $a = 3, 2, 1$ parameterize each gauge group of $SU(3)_C \times SU(2)_L \times U(1)_Y$. For each gauge group G_a , we introduce an adjoint chiral multiplet A_a which contains an R -partner χ_a of the gaugino λ_a . Since the gaugino λ_a has R -charge $+1$, its partner χ_a has to have R -charge -1 so that the adjoint chiral superfield A_a has $R = 0$. Accordingly, we have a Dirac type mass term

$$\mathcal{L}_{\text{gaugino}} = - \sum_{a=1,2,3} m_a \lambda_a \chi_a + \text{H.c.} . \quad (2)$$

The Dirac gaugino mass term can be generated through the “supersoft” operator [2]

$$\mathcal{L}_{\text{supersoft}} = \sum_{a=1,2,3} \int d^2\theta \sqrt{2} \frac{W'^a W_\alpha^a A_a}{\Lambda_D} + \text{H.c.} , \quad (3)$$

where Λ_D is a messenger scale, W_α^a is the MSSM gauge field strength, W'_α is a hidden-sector gauge field strength which acquires a nonzero D -term, $\langle W'_\alpha \rangle = \theta_\alpha \langle D' \rangle$ so that $m_a = \langle D' \rangle / \Lambda_D$. In the present work, we assume that a suitable size of masses are generated although we do not elucidate the hidden sector dynamics.

Next, we turn to the Higgs sector. Following the Ref. [4], we consider the superpotential

$$W_{\text{Higgs}} = X_0 (f + \lambda H_u H_d) + \mu_1 X_d H_u + \mu_2 X_u H_d . \quad (4)$$

The gauge and R -charge assignments are shown in Table. I. The superfields X_d and X_u with R -charge 2 are the mirror partners of the MSSM Higgs fields H_u and H_d , with the supersymmetric masses μ_1 and μ_2 , respectively: The $SU(2) \times U(1)$ singlet field X_0 also has R -charge 2. The dimension two parameter f is the source of the visible SUSY breaking while the dimensionless coupling λ plays important roles not only for triggering the EWSB but also for generating the quartic coupling of the MSSM Higgs scalars.

	X_0	X_u	X_d	H_u	H_d	L	E	N
$SU(2)_L$	1	2	2	2	2	2	1	1
$U(1)_Y$	0	+1/2	-1/2	+1/2	-1/2	-1/2	+1	0
$U(1)_R$	+2	+2	+2	0	0	+1	+1	-1

TABLE I. The charge assignment of the Higgs and lepton fields under the EW symmetry and $U(1)_R$ symmetry. The R -charges are those for left-handed chiral superfields. The MSSM Higgs doublets are R -neutral while their R -partners and the singlet X_0 have R -charge +2. All the quarks and leptons have R -charge +1, except that the right-handed neutrinos have R -charge -1. The implications of this flipped assignment will be discussed in Sec. III. Note also that the adjoint chiral multiplets A_a are R -neutral so that their fermionic components have R -charge -1.

With the R -symmetric superpotential (4), supersymmetry is spontaneously broken in accordance with the general argument³. Moreover, such visible SUSY breaking triggers the correct EWSB if the coupling λ is sufficiently large. A possible origin of the dimensionful parameters f and $\mu_{1,2}$ was suggested in Ref. [4]. Here we just assume that the scale of these parameters is around the weak scale; we will take $\mu_{u,d} = 300$ GeV in our analysis in the next section.

We note that the adjoint chiral multiplets $A_{a=2,1}$ of $SU(2) \times U(1)$ can have a $U(1)_R$ -symmetric superpotential interactions with the Higgs fields $H_{u,d}$ and $X_{d,u}$. For simplicity, however, we do not include them in our analysis. More importantly, we do not include the following mixing term between the singlet X_0 and the $U(1)_Y$ “adjoint” multiplet A_1 .

$$W_{\text{mix}} = \mu_0 X_0 A_1 . \quad (5)$$

This term is very dangerous since it could cancel, if present, the linear term in the superpotential (4). The absence of such term can be justified if $U(1)_Y$ is embedded into a simple gauge group at high energy. See Ref. [2] for a similar discussion about the absence of the kinetic mixing of the hidden $U(1)'$ and $U(1)_Y$.

B. Minimal $U(1)_R$ Symmetry Breaking

Next we discuss $U(1)_R$ symmetry breaking. Once the model is coupled to supergravity, the $U(1)_R$ symmetry is necessarily broken in order that the cosmological constant can be adjusted to zero. Our basic assumption here is that such breaking of R symmetry is mediated to the visible sector in a minimal way; under such assumption of “minimal R -breaking mediation”, the interaction Lagrangian is given by⁴

$$\mathcal{L}_{\text{eff}} = \int d^2\theta \phi^3 W + \text{H.c.} , \quad \phi = 1 + \theta^2 m_{3/2} , \quad (6)$$

³ There exists no supersymmetric vacuum when the number of the fields with R -charge 2 is larger than the number of the fields with R -charge 0. See Ref. [14] and references therein.

⁴ Here we took a modified gauge fixing of Ref. [15].

where ϕ is the so-called conformal compensator and $m_{3/2}$ is the mass of the gravitino. The conformal compensator ϕ can be absorbed by rescaling the chiral superfields, $\phi\Phi_i \rightarrow \Phi_i$, where Φ_i represents all the chiral superfields in the theory. Classically such rescaling has no effect on cubic terms in superpotential, whereas terms of dimensions less than four are affected. In the present case, the interaction Lagrangian takes the form

$$\mathcal{L}_{\text{eff}} = \int d^2\theta W + m_{3/2}G(\Phi_i) + \text{H.c.} , \quad (7)$$

$$G(\Phi_i) \equiv 2fX_0 + \mu_1X_dH_u + \mu_2X_uH_d , \quad (8)$$

where the fields in $G(\Phi_i)$ are the scalar components of the corresponding superfields $\Phi_i = H_u, H_d, X_u, X_d$ or X_0 . Terms proportional to the gravitino mass represent R -breaking interactions, and are small if the gravitino mass is small; we assume throughout the present paper that the gravitino mass is much smaller than the weak scale. As we will see shortly, however, these interactions induce a slight shift of the vacuum since they contain a tadpole term of X_0 and also those of $X_{u,d}$ after the EWSB. Such shift of the VEVs will play important roles when we discuss the generation of Dirac neutrino masses (in §III) and properties of the pseudo goldstino (in §IV).

C. The EWSB Vacuum

We now analyze the scalar potential in order to find the shift of the VEVs induced by the R -breaking effects. Here we assume that the Kähler potential is canonical. The scalar potential involving the neutral Higgs fields $H_{u,d}^0$ and their mirror partners $X_{u,d}^0$ and the singlet X_0 is given by

$$V(\Phi_i, \Phi_i^\dagger) = V_0(\Phi_i, \Phi_i^\dagger) - \{m_{3/2}G(\Phi_i) + \text{H.c.}\} , \quad (9)$$

with the R -symmetric part V_0 given by

$$\begin{aligned} V_0 = & |f - \lambda H_u^0 H_d^0|^2 + |\lambda X_0 H_d^0 - \mu_1 X_d^0|^2 + |\lambda X_0 H_u^0 + \mu_2 X_u^0|^2 \\ & + \frac{1}{8}g^2 \left(|H_u^0|^2 - |H_d^0|^2 + |X_u^0|^2 - |X_d^0|^2 \right)^2 \\ & + \bar{m}_{H_u}^2 |H_u^0|^2 + \bar{m}_{H_d}^2 |H_d^0|^2 - (bH_u^0 H_d^0 + \text{h.c.}) \\ & + m_{X_0}^2 |X_0|^2 + m_{X_u}^2 |X_u^0|^2 + m_{X_d}^2 |X_d^0|^2 . \end{aligned} \quad (10)$$

The first and the second lines represent the F -term and the D -term potentials, respectively.⁵ The coupling constant g^2 is $g_1^2 + g_2^2$ where g_1 and g_2 are the gauge coupling constants of $U(1)$ and $SU(2)$ symmetries, respectively. The remaining are soft SUSY breaking terms except that the masses of the MSSM Higgs fields, $\bar{m}_{H_u}^2$ and $\bar{m}_{H_d}^2$, are the sums of soft masses and the supersymmetric masses. We suppose that these soft terms are induced radiatively from

⁵ As noted in Ref. [2] the $SU(2) \times U(1)$ D -terms are absent in the supersoft limit. If the adjoint scalars get soft scalar masses through R -invariant mediation of SUSY breaking, then D terms does not decouple completely. Our results in the present paper are not affected whether the D -terms decouple or not.

the supersoft operators, or mediated from the hidden sector in an R -invariant way. We also assume that all dimensionful parameters are around or just above the weak scale whereas the gravitino mass is much smaller. In the following, we consider the case in which $m_{X_0}^2$ and $m_{X_{u,d}}^2$ are positive so that X_0 and $X_{u,d}$ do not develop VEVs in the R -symmetric limit.

The vacuum with EWSB can be found by solving the stationary conditions

$$0 = \frac{\partial V}{\partial \Phi_i} = \frac{\partial V_0}{\partial \Phi_i} - m_{3/2} \frac{\partial G(\Phi_i)}{\partial \Phi_i} . \quad (11)$$

We solve these equations perturbatively with respect to a small gravitino mass. In the R -symmetric limit, *i.e.*, $m_{3/2} \rightarrow 0$, the solution was given in Ref. [4], which we denote by $\langle \Phi_i \rangle_0$. Note that the unbroken $U(1)_R$ symmetry implies $\langle X_0 \rangle_0 = \langle X_u \rangle_0 = \langle X_d \rangle_0 = 0$.

Now, with the R -breaking terms, the solution to Eq. (11) takes the form

$$\langle \Phi_i \rangle = \langle \Phi_i \rangle_0 + \kappa_i m_{3/2} + \mathcal{O}(m_{3/2}^2) , \quad (12)$$

where the coefficient κ_i is found to be

$$\kappa_i = M_{ij}^{-1} \left. \frac{\partial G}{\partial \Phi_j} \right|_{\Phi=\langle \Phi \rangle_0} , \quad M_{ij} \equiv \left. \frac{\partial^2 V_0}{\partial \Phi_i \partial \Phi_j} \right|_{\Phi=\langle \Phi \rangle_0} . \quad (13)$$

Explicitly we obtain, for the R -charged Higgs fields,

$$\langle X_0 \rangle = \frac{2f}{m_{X_0}^2 + \lambda^2 v^2} m_{3/2} , \quad (14)$$

$$\langle X_d^0 \rangle = \frac{\mu_1 v \sin \beta}{m_{X_d}^2 + \mu_1^2 + (m_Z^2/2) \cos 2\beta} m_{3/2} , \quad (15)$$

$$\langle X_u^0 \rangle = \frac{-\mu_2 v \cos \beta}{m_{X_u}^2 + \mu_2^2 - (m_Z^2/2) \cos 2\beta} m_{3/2} , \quad (16)$$

where $m_Z^2 = (g_1^2 + g_2^2) v^2/2$ is the Z boson mass.⁶ We see that the VEV of X_0 is proportional to f that represents the scale of the visible SUSY breaking, while the VEVs of $X_{u,d}$ are proportional to their supersymmetric masses and their partner's VEVs. These proportionality can be understood if one notices that replacing the Higgs fields by their VEVs in Eq. (8) generates the tadpole term for each of $X_{0,u,d}$,

$$G(X_0, X_u^0, X_d^0) = 2f X_0 + \mu_1 \langle H_u^0 \rangle X_d^0 - \mu_2 \langle H_d^0 \rangle X_u^0 . \quad (17)$$

Of course, the VEVs of these R -charged fields are all proportional to the gravitino mass⁷ which parameterizes the $U(1)_R$ -symmetry breaking.

⁶ Note that the VEVs of the MSSM Higgs fields also receive $\mathcal{O}(m_{3/2})$ shifts, which are negligibly small for a small gravitino mass.

⁷ Precisely speaking, the VEVs $\langle X_{0,u,d} \rangle$ have R -charge $+2$ and should be proportional to $m_{3/2}^*$ since one can regard the gravitino mass as having R -charge -2 in the sense of spurion analysis.

III. NEUTRINO MASS AND PROTON DECAY IN FLIPPED $U(1)_R$ MODEL

In this section, we show that the neutrinos can acquire tiny masses via the supergravity-induced effect of $U(1)_R$ -symmetry breaking. We introduce three right-handed neutrino multiplets N_i^c ($i = 1, 2, 3$) and assign the R -charge -1 to them as shown in Table. I. That is, we flip the R -charge of the right-handed neutrinos. Under this flipped assignment, we first argue that the neutrinos can be Dirac particles, and that the smallness of neutrino masses are related to the smallness of $U(1)_R$ -symmetry breaking parameterized by the gravitino mass. Then, we discuss the proton decay induced by a new type of dimension five operator involving the right-handed neutrinos.

A. Neutrino Mass in Flipped $U(1)_R$ Model

Under our flipped $U(1)_R$ assignment, the right-handed neutrino can have a Yukawa-type interaction with the mirror Higgs field X_u , instead of the MSSM Higgs H_u . At the renormalizable level the matter superpotential is given by

$$W = y_U^{ij} Q_i U_j^c H_u + y_D^{ij} Q_i D_j^c H_d + y_E^{ij} E_i^c L_j H_d + y_N^{ij} N_i^c L_j X_u , \quad (18)$$

where $y_{U,D,E,N}$ are 3×3 Yukawa coupling matrices.

The Majorana mass term for the right-handed neutrinos is forbidden by our R -charge assignment. The symmetry allows the Weinberg operator

$$W_{\Delta L=2} = - \frac{C_{\Delta L=2}^{ij}}{\Lambda_{\text{cutoff}}} (L_i H_u) (L_j H_u) , \quad (19)$$

which gives a sub-dominant contribution to the neutrino mass matrix if we assume the cutoff Λ_{cutoff} to be larger than 10^{12} – 10^{13} GeV. Therefore the neutrinos are almost Dirac particles in our model.

By replacing the mirror Higgs field with its VEV (16), we obtain the Dirac mass term for neutrinos as

$$\mathcal{L}_{\nu \text{ mass}} = - m_\nu \nu_R^c \nu_L + \text{H.c.} , \quad m_\nu^{ij} = y_N^{ij} \langle X_u^0 \rangle . \quad (20)$$

Since we are most interested in the heaviest neutrino, we will suppress the flavor indices hereafter and denote the heaviest eigenvalue and the corresponding Yukawa coupling by m_ν and y_ν , respectively. With this understanding, we write

$$m_\nu = y_\nu \langle X_u^0 \rangle = y_\nu m_{3/2} \left[\frac{-\mu_2 v \cos \beta}{m_{X_u}^2 + \mu_2^2 - (m_Z^2/2) \cos 2\beta} \right] . \quad (21)$$

A notable feature in the present model is that the size of the neutrino mass is set by that of the gravitino mass $m_{3/2}$ when all other dimensionful parameters are of the same order. Accordingly the neutrino Yukawa coupling can be of $\mathcal{O}(1)$ if the gravitino is as light as 10 eV. In fact, as we mentioned in Eq. (1), the mass of the gravitino is constrained cosmologically;

$$m_{3/2} \lesssim 16 \text{ eV} , \quad 10 \text{ keV} \lesssim m_{3/2} .$$

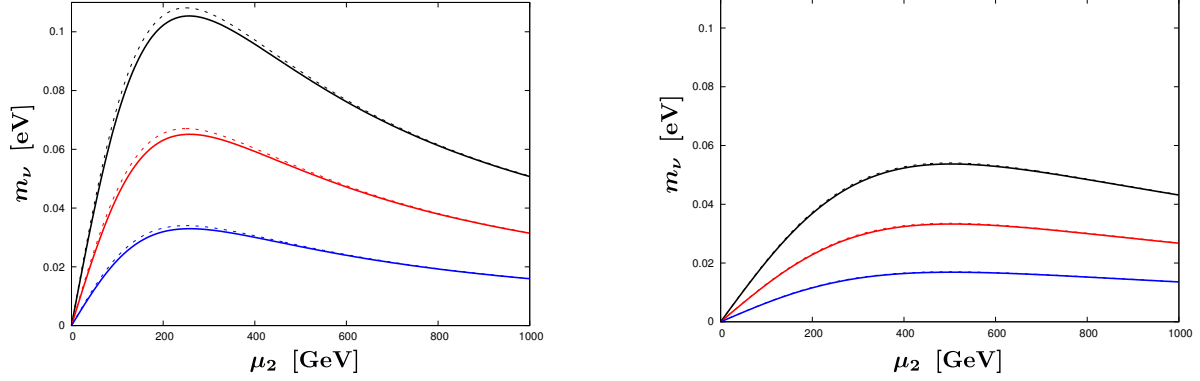


FIG. 1. Left: The heaviest neutrino mass as a function of the Higgsino mass parameter μ_2 , with $m_{X_d} = 250$ GeV (Left) and 500 GeV (Right). In each figure, the solid and dotted lines correspond to the case with and without the D -term contribution in the denominator in Eq. (21). Three lines correspond to $\tan \beta = 3$ (top), $\tan \beta = 5$ (middle), and $\tan \beta = 10$ (bottom), respectively.

In the case of lighter gravitino, a tiny neutrino mass can be obtained with a large Yukawa coupling. Even in the case of heavier gravitino, the neutrino Yukawa can be comparable to the electron Yukawa.

Fig. 1 shows how the generated neutrino mass depends on the Higgs mass parameters, for a fixed value of $m_{3/2} = 10$ eV and $y_\nu = 0.1$. We see that the neutrino mass is maximized for $\mu_2 \sim m_{X_d}$. We also note that the D -term contribution in the denominator in Eq. (21) can safely be neglected when the Higgs mass parameters are larger than m_Z . Accordingly the neutrino mass is proportional to $\cos \beta$ instead of $\sin \beta$.

Some remarks are in order. First, our $U(1)_R$ -charge assignment does not induce the mixed anomalies with the SM gauge groups since the right-handed neutrinos are SM singlets. $U(1)_R$ anomalies can be canceled by introducing other singlets with positive $U(1)_R$ -charges, without changing our results below. Second, our model can be regarded as a realization of “neutrinophilic Higgs” idea: the original idea was proposed in non-SUSY context in Refs. [16, 17] in which a softly broken Z_2 symmetry is the source of the tiny VEV. See also Ref. [18] for a softly broken global $U(1)$ case. Quantum stability was discussed in Ref.[19, 20]. The point in our model is that the size of the “neutrinophilic” Higgs VEV is related to that of the gravitino mass, and we can address its implications in cosmology and Higgs phenomenology, in addition to the proton stability as we shall discuss shortly. We also note that many other topics were discussed in literature, such as muon $g - 2$ and lepton flavor violations [17, 18], low-scale leptogenesis [21], dark matter and cosmology [22], dark energy [23], and supernova neutrinos [24]: In particular, Ref. [25] claimed that the size of neutrino Yukawa couplings are severely constrained from the observations of supernova neutrinos as well as CMB radiations, if the neutrinophilic Higgs scalar is extremely light. This constraint does not apply here since the corresponding scalar is heavy enough.

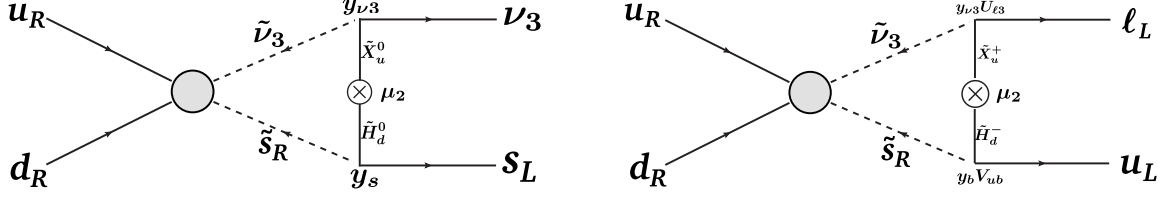


FIG. 2. Left: A diagram with neutral higgsino dressing gives the dominant mode $p \rightarrow \bar{\nu} + K^+$. Right: Charged higgsino dressing diagrams give sub-dominant modes $p \rightarrow \ell^+ + \pi^0$ ($\ell = e, \mu$). For the latter, the diagrams with a sbottom \tilde{b}_R in the loop also give a comparable contribution.

B. Proton Decay in Flipped $U(1)_R$ Model

Originally, a nice feature of $U(1)_R$ symmetry is that it explains proton stability naturally: the dangerous baryon and lepton number violating operators are forbidden if we assign R -charge $+1$ to all the quark and lepton superfields and R -charge 0 to the Higgses. This nice property is not modified when we introduce the R -partner Higgses $X_{0,u,d}$ with $R = 2$ and the R -partner gauginos with $R = 0$. In our assignment, however, we flip the R -charge of the right-handed neutrinos N^c to be -1 , which may spoil the proton stability.

An operator analysis presented in Appendix A shows that there is a unique dimension-five operator that can lead to the proton decay,

$$W_{5R} = - \frac{C_5^{ijkl}}{\Lambda_{\text{cutoff}}} U_i^c D_j^c D_k^c N_\ell^c, \quad (22)$$

where Λ_{cutoff} is a cutoff scale and C_5^{ijkl} is a dimensionless coefficient. In this subsection we will give a rough estimate of the proton lifetime. To simplify the expressions, we will work in a flavor basis in which the Yukawa couplings of up-type quarks and charged leptons are diagonal; we also define C_5^{ijkl} to be the coefficient of the ℓ -th mass eigenstate of neutrinos.

With the dimension five operator (22), the proton decay occurs via the processes depicted in Fig. 2, the one with neutral Higgsino dressing and the other with charged Higgsino dressing. Remarkably, diagrams involving the top Yukawa coupling are absent because the $\mu H_u H_d$ term is absent in the $U(1)_R$ -symmetric limit, or is extremely suppressed as is proportional to the tiny VEV (14). Instead, we have potentially large contributions involving a large neutrino Yukawa coupling y_ν of the heaviest neutrino ν_3 .

The partial rate for a proton decaying into a meson M and a lepton ℓ takes the form⁸

$$\Gamma(p \rightarrow M + \ell) = \frac{m_p}{32\pi} \left(1 - \frac{m_M^2}{m_p^2}\right)^2 \frac{|\alpha_p|^2}{f_M^2} \left| \mathcal{A}(p \rightarrow M\ell) \right|^2. \quad (23)$$

Here M_M and f_M are the mass and the decay constant of the meson M , respectively. The dimensionful constant α_p is defined through $\langle 0 | \epsilon_{abc} (d_R^a u_R^b) u_L^c | 0 \rangle = \alpha_p N_L$ where a, b, c are color indices and N_L is the wavefunction of left-handed proton. We use the value

⁸ For long distance effects of RRRR-type operator, see Ref. [26].

$\alpha_p = -0.015 \text{ GeV}^3$ [27] in our calculation. The amplitude \mathcal{A} corresponding to each diagram in Fig. 2 is given respectively by

$$\mathcal{A}(p \rightarrow K^+ \bar{\nu}_3) = \frac{C_5^{1123}}{\Lambda_{\text{cutoff}}} \frac{y_{\nu 3} y_s \mu_2}{16\pi^2 m_{\text{soft}}^2}, \quad (24)$$

$$\mathcal{A}(p \rightarrow \pi^0 \ell^+) = \sum_{k=s,b} \frac{C_5^{11k3}}{\Lambda_{\text{cutoff}}} \frac{U_{\ell 3} y_{\nu 3} V_{uk} y_k \mu_2}{16\pi^2 m_{\text{soft}}^2}, \quad (25)$$

where $U_{\ell 3}$ and V_{uk} are mixing matrix elements, and m_{soft} is a typical mass scale of the particles propagating in the loop. We see that the latter diagrams with charged higgsino dressing are suppressed by the CKM matrix element $\varepsilon = V_{us} = 0.22$, because $y_s \sim y_b \varepsilon^2$ and $y_b V_{ub} \sim y_s V_{us} \sim y_b \varepsilon^3$. Therefore the dominant decay mode is $p \rightarrow K^+ + \bar{\nu}$. Using the simplified expression of the neutrino Yukawa⁹

$$y_\nu \approx \frac{m_\nu}{m_{3/2}} \left(\frac{m_{X_u}^2 + \mu_2^2}{-\mu_2 v \cos \beta} \right), \quad (26)$$

we find the partial decay width into a kaon and an anti-neutrino is given by

$$\Gamma(p \rightarrow K^+ \bar{\nu}) \approx \frac{m_p}{32\pi} \left(1 - \frac{m_{K^+}^2}{m_p^2} \right)^2 \frac{\alpha_p^2}{f_K^2} \left(\frac{C_5}{16\pi^2 m_{3/2} \Lambda_{\text{cutoff}}} \right)^2 \frac{m_s^2 m_\nu^2}{v^4 \cos^4 \beta} \left(\frac{\mu_2^2 + m_{X_u}^2}{m_{\text{soft}}^2} \right)^2, \quad (27)$$

where we denote C_5^{1123} simply by C_5 .

Numerically the lifetime of proton can be estimated as

$$\tau_p = \frac{5.65 \times 10^{28} \text{ yr}}{C_5^2} \left(\frac{\Lambda_{\text{cutoff}}}{2.4 \times 10^{18} \text{ GeV}} \right)^2 \left(\frac{m_{3/2}}{10 \text{ eV}} \right)^2 \left(\frac{\cos^2 \beta}{0.1} \right)^2, \quad (28)$$

where we took the following values of parameters:

$$\begin{aligned} m_{\text{soft}} &= 1.5 \text{ TeV}, & \mu_2 &= m_{X_u} = 300 \text{ GeV}, & \tan \beta &= 3, \\ m_p &= 1.0 \text{ GeV}, & \alpha_p &= -0.015 \text{ GeV}^3, \\ m_{K^+} &= 0.5 \text{ GeV}, & f_K &= 0.13 \text{ GeV}, \\ m_s &= 0.1 \text{ GeV}, & m_\nu &= 0.1 \text{ eV}. \end{aligned} \quad (29)$$

The present lower bound on the proton lifetime is $2.3 \times 10^{33} \text{ yr}$ (90%CL) obtained for $p \rightarrow K^+ + \bar{\nu}$ mode [28]. By comparing the bound with the result, we obtain the constraint¹⁰

$$|C_5| \lesssim 4.96 \times 10^{-3} \left(\frac{\Lambda_{\text{cutoff}}}{2.4 \times 10^{18} \text{ GeV}} \right) \left(\frac{m_{3/2}}{10 \text{ eV}} \right) \left(\frac{\cos^2 \beta}{0.1} \right). \quad (30)$$

Fig. 3 shows the constraint from proton stability as a function of the gravitino mass. The constraint is most severe for a small gravitino mass: If the gravitino is as light as $\mathcal{O}(\text{eV})$,

⁹ Here we have neglected the mild $\tan \beta$ dependence in the denominator in Eq. (21) to simplify the expression. We used Eq. (21) in our numerical calculation.

¹⁰ This applies only for $\Lambda_{\text{cutoff}} \gtrsim 10^{13} \text{ GeV}$; otherwise, the operators (19) can not be neglected anymore.

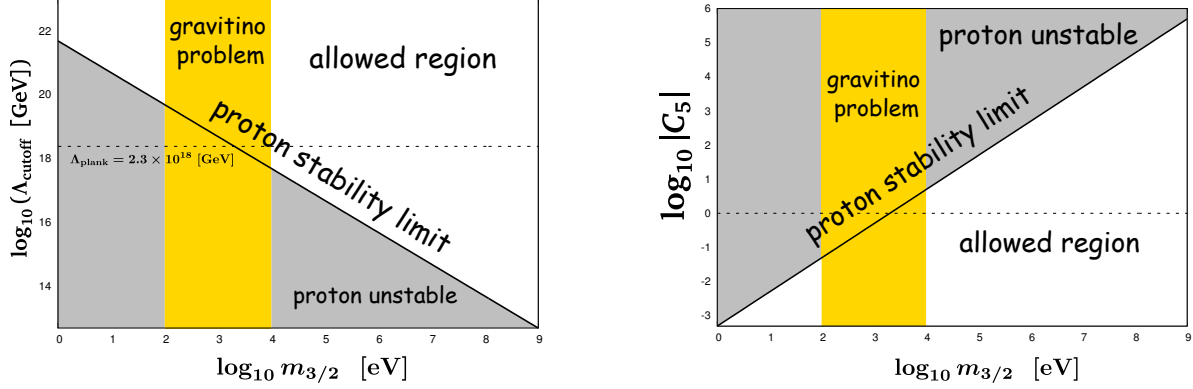


FIG. 3. Constraints from proton stability as a function of the gravitino mass. Left: the constraint on $\log_{10} \Lambda_{\text{cutoff}}$ for the fixed coupling $C_5 = 1$. Right: the constraint on $\log_{10} |C_5|$ for $\Lambda_{\text{cutoff}} = M_{\text{pl}}$. We took $\tan \beta = 3$, $\mu_2 = 300$ GeV, $m_{X_u} = 300$ GeV, and $m_{\tilde{q}} = 1.5$ TeV. The figures correspond to the neutrino mass $m_\nu = 0.1$ eV. The gray regions are excluded by the proton decay while the yellow bands are disfavored cosmologically.

the coefficient C_5 need to be a few order of magnitudes smaller than 1 even when the cutoff Λ_{cutoff} is equal to the reduced Planck scale M_{pl} . The constraint becomes milder for a larger gravitino mass. This can be understood from Eq. (21): For a fixed neutrino mass, a larger gravitino mass implies a smaller neutrino Yukawa coupling, suppressing the proton decay. We thus find that the constraint can be satisfied for the gravitino mass larger than 10 keV, which is also viable cosmologically as in seen in Eq. (1).

Notice, however, that the constraint becomes severe for a large $\tan \beta$: The lifetime is proportional to $\cos^4 \beta$ up to a possible mild dependence in the denominator in Eq. (21). This is because the X_u^0 tadpole is proportional to the $\langle H_d^0 \rangle$ and because the decay amplitude involves a down-type Yukawa coupling. Consequently, for $\tan \beta = 10$, the proton stability against the dimension five operator (22) requires that the gravitino mass should be larger than $\mathcal{O}(100 \text{ keV})$. In this case, the neutrino Yukawa coupling y_ν is comparable to the electron Yukawa.

The above results should be compared with the MSSM case in which the proton decay is induced by the dimension five operators

$$\Delta W_{\text{MSSM}} = -\frac{C_L}{2\Lambda_{\text{cutoff}}} QQQQL - \frac{C_R}{\Lambda_{\text{cutoff}}} U^c U^c D^c E^c. \quad (31)$$

In the MSSM, the dominant contribution comes from the LLLL operator, and the coefficient should be very suppressed, $|C_L| \lesssim 10^{-8}$ for $\Lambda_{\text{cutoff}} = M_{\text{pl}}$. Moreover, the RRRR operator should also be suppressed because it can involve the top Yukawa coupling when dressed with a charged Higgsino loop. In our case, the LLLL operator is absent due to the R -symmetry, and with our RRRR operator, chargino loop diagrams in Fig. 2 involve a down-type Yukawa coupling, y_s or $y_b V_{cb}$, instead of top Yukawa. As a consequence, the RRRR contribution in the flipped $U(1)_R$ model is a couple of order smaller than that in the MSSM.

In this way, our flipped $U(1)_R$ assignment can explain the smallness of neutrino masses with relatively large Yukawa couplings and without spoiling the proton stability so much.

The latter is certainly true for small $\tan\beta$: for larger $\tan\beta$, on the other hand, there is a tension between our mechanism for neutrino masses and the proton stability.

We note that such tension can be relaxed if we combine the idea of flavor symmetries with the present model. In Appendix B, we give an illustrative example of $U(1)$ flavor symmetry, along the line of Ref. [29], and show that a proper $U(1)$ charge assignment for generating the Yukawa hierarchy can guarantee the proton stability as well.

IV. MASS AND MIXING OF PSEUDO GOLDSTINO

We now turn to another effect of R -symmetry breaking in our $U(1)_R$ -symmetric model. When supersymmetry is broken in two independent sectors, there appear two Goldstine fermions [11]. After coupling to supergravity, one linear combination of these goldstinos becomes the longitudinal components of the massive gravitino, and the other is a pseudo goldstino state, which we denote by ζ . The mass of such pseudo goldstino has been studied in literature [6, 11, 12].

In our case, supersymmetry is broken in the visible sector as well as in a hidden sector. For our purpose, we do not need to specify its precise form of the hidden sector, but we just assume that the supersymmetry breaking is hierarchical; the SUSY breaking scale in the hidden sector is much larger than that in the visible sector. We then expect that the physical pseudo goldstino state ζ resides dominantly in the visible sector. In the limit $\lambda \rightarrow 0$, in which the EWSB is switched off, the would-be goldstino in the visible sector is the singlino \tilde{X}_0 , the fermionic component of the singlet X_0 , since the visible SUSY breaking is triggered by its linear term in the superpotential (4). Therefore we first discuss the mass of X_0 fermion, and then calculate the full neutralino mass matrix and its smallest eigenvalue and the eigenvector.

A. Mass of Singlet Fermion from $U(1)_R$ Breaking

The mass term of the singlet fermion \tilde{X}_0 can be generated from the contact term in the Kähler potential of X_0 . For definiteness, let us consider the following Kähler potential

$$K(X_0, X_0^\dagger) = X_0 X_0^\dagger - \frac{1}{4\Lambda_0^2} (X_0 X_0^\dagger)^2, \quad (32)$$

where Λ_0 is the cutoff scale at which the contact term is generated. The relevant terms in the Lagrangian containing the X_0 supermultiplet are given by

$$\begin{aligned} \mathcal{L}_{X_0} = & K_{X_0^\dagger X_0} F_{X_0}^\dagger F_{X_0} - m_{X_0}^2 X_0^\dagger X_0 - \lambda^2 (|H_u|^2 + |H_d|^2) X_0^\dagger X_0 \\ & + \left\{ F_{X_0} \frac{\partial W}{\partial X_0} - \frac{1}{2} m_{\tilde{X}_0} \tilde{X}_0 \tilde{X}_0 + \text{H.c.} \right\} + \{ 2f m_{3/2} X_0 + \text{H.c.} \}, \end{aligned} \quad (33)$$

where the (Majorana) mass of the would-be goldstino \tilde{X}_0 is found to be

$$m_{\tilde{X}_0} = -\frac{1}{\Lambda_0^2} \left\langle \int d^2\bar{\theta} \frac{1}{2} X_0^\dagger X_0^\dagger \right\rangle = -\frac{\langle F_{X_0} \rangle^\dagger}{\Lambda_0^2} \langle X_0 \rangle^\dagger. \quad (34)$$

Note that this expectation value can be non-vanishing only in the presence of the $U(1)_R$ breaking, the last term in the Lagrangian (33). To compute it, we use the equations of motion to get

$$\langle F_{X_0}^\dagger \rangle = -K_{X_0^\dagger X_0}^{-1} \frac{\partial W}{\partial X_0} \approx -(f - \lambda \langle H_u^0 \rangle \langle H_d^0 \rangle) , \quad (35)$$

$$\langle X_0 \rangle = \frac{2f}{m_{X_0}^2 + \lambda^2 v^2 + \delta m_{X_0}^2} m_{3/2} . \quad (36)$$

Here in the first equation, we have used the fact that the Kähler metric is almost canonical since $|\langle X_0 \rangle| \ll \Lambda_0$. In the second equation, we note that the previous result (14) is slightly modified by the soft scalar mass $\delta m_{X_0}^2$ due to the contact term in Eq. (32),

$$\delta m_{X_0}^2 \equiv \frac{|\langle F_{X_0} \rangle|^2}{\Lambda_0^2} \approx \frac{1}{\Lambda_0^2} |f - \lambda v^2 \sin \beta \cos \beta|^2 . \quad (37)$$

By plugging the VEVs of X_0 and F_{X_0} into the expression (34), we thus find that the X_0 fermion mass is proportional to the gravitino mass as

$$m_{\tilde{X}_0} \approx 2\mathcal{Z}_1 m_{3/2} , \quad \mathcal{Z}_1 \equiv \frac{f(f - \lambda v^2 \sin \beta \cos \beta)}{\Lambda_0^2 (m_{X_0}^2 + \lambda^2 v^2 + \delta m_{X_0}^2)} . \quad (38)$$

The result (38) is consistent with the general assertion [11] that the pseudo goldstino mass is twice the gravitino mass in the “sequestered” limit. This can be seen as follows: If the soft mass of X_0 is zero, $m_{X_0}^0 \rightarrow 0$, and if the Higgs sector decouples from the visible SUSY breaking, $\lambda \rightarrow 0$, we have $\delta m_{X_0}^2 \rightarrow f^2/\Lambda_0^2$, and hence,

$$\mathcal{Z}_1 \longrightarrow 1 . \quad (39)$$

Otherwise, the pseudo goldstino mass is suppressed by a factor \mathcal{Z}_1 , which is roughly of order f/Λ_0^2 , in accordance with the general argument [6]. For instance, for a moderate choice of parameters $f \sim (10^3 \text{ GeV})^2$ and $\Lambda_0 \sim 10^7 \text{ GeV}$, the mass of the pseudo goldstino is smaller than $\mathcal{O}(10 \text{ eV})$ even if the gravitino is as heavy as $m_{3/2} \sim \text{GeV}$.

B. Neutralino Mass Matrix with Dirac Gaugino

Now we consider the neutralino mass matrix which incorporates the pseudo goldstino. In this subsection, let us denote by χ_B and χ_W the Dirac partner of the bino \tilde{B} and the wino \tilde{W} , respectively. The Dirac mass terms for $SU(2) \times U(1)$ gauginos are

$$\mathcal{L}_{\text{gaugino}} = -m_{\tilde{B}} \chi_B \tilde{B} - m_{\tilde{W}} \chi_W \tilde{W} + \text{H.c.} . \quad (40)$$

Then, in a basis given by

$$\vec{\Psi}^T = \left(\begin{array}{cccc|cccc} \chi_B & \chi_{W^3} & \tilde{W}^3 & \tilde{B} & \tilde{H}_d^0 & \tilde{H}_u^0 & \tilde{X}_d^0 & \tilde{X}_u^0 \end{array} \middle| \tilde{X}_0 \right) , \quad (41)$$

the neutralino mass matrix takes the form

$$\mathcal{M}_{\tilde{N}} = \left(\begin{array}{cc|cc|cc|cc|c} 0 & 0 & 0 & m_{\tilde{B}} & 0 & 0 & 0 & 0 & 0 \\ 0 & 0 & m_{\tilde{W}} & 0 & 0 & 0 & 0 & 0 & 0 \\ \hline 0 & m_{\tilde{W}} & 0 & 0 & am_Z & bm_Z & a'm_V & b'm_V & 0 \\ m_{\tilde{B}} & 0 & 0 & 0 & cm_Z & dm_Z & c'm_V & d'm_V & 0 \\ \hline 0 & 0 & am_Z & cm_Z & 0 & -\lambda v_0 & 0 & +\mu_2 & -\lambda v_u \\ 0 & 0 & bm_Z & dm_Z & -\lambda v_0 & 0 & -\mu_1 & 0 & -\lambda v_d \\ \hline 0 & 0 & a'm_V & c'm_V & 0 & -\mu_1 & 0 & 0 & 0 \\ 0 & 0 & b'm_V & d'm_V & +\mu_2 & 0 & 0 & 0 & 0 \\ \hline 0 & 0 & 0 & 0 & -\lambda v_u & -\lambda v_d & 0 & 0 & m_{\tilde{\chi}_0} \end{array} \right), \quad (42)$$

where $m_{\tilde{\chi}_0} = 2m_{3/2}\mathcal{Z}_1$, and we have defined $\langle X^0 \rangle = v_0$ and

$$m_V^2 \equiv \frac{g_1^2 + g_2^2}{2} \left(|\langle X_d^0 \rangle|^2 + |\langle X_u^0 \rangle|^2 \right), \quad \tan \gamma \equiv \frac{\langle X_u^0 \rangle}{\langle X_d^0 \rangle}. \quad (43)$$

We have also used the following abbreviations: with θ_W being the weak mixing angle,

$$U \equiv \begin{pmatrix} a & b \\ c & d \end{pmatrix} \equiv \begin{pmatrix} c_\beta c_W & -s_\beta c_W \\ -c_\beta s_W & s_\beta s_W \end{pmatrix} = \begin{pmatrix} \cos \theta_W \\ -\sin \theta_W \end{pmatrix} \begin{pmatrix} \cos \beta & -\sin \beta \end{pmatrix}, \quad (44)$$

where $\langle H_d^0 \rangle = v \cos \beta$ and $\langle H_u^0 \rangle = v \sin \beta$: the primed quantities are obtained by replacing β with the angle γ defined above.

The mass eigenvalue Ω_D can be found by solving the characteristic equation

$$0 = \det(\mathcal{M}_{\tilde{N}} - \Omega_D \hat{1}) = \det \mathcal{M}_{\tilde{N}} - \Omega_D \det \mathcal{M}_{\tilde{N}} \text{Tr}[\mathcal{M}_{\tilde{N}}^{-1}] + \dots, \quad (45)$$

where the dots represents terms higher order in Ω_D . The first and the second terms are calculated to be

$$\det \mathcal{M}_{\tilde{N}} = 2\mathcal{Z}_1 m_{3/2} \mu_1^2 \mu_2^2 m_{\tilde{B}}^2 m_{\tilde{W}}^2, \quad (46)$$

$$\det \mathcal{M}_{\tilde{N}} \text{Tr}[\mathcal{M}_{\tilde{N}}^{-1}] = m_{\tilde{B}}^2 m_{\tilde{W}}^2 [\mu_1^2 \mu_2^2 + \lambda^2 v^2 (\mu_1^2 \sin^2 \beta + \mu_2^2 \cos^2 \beta)], \quad (47)$$

up to $\mathcal{O}(m_{3/2}^2)$ and $\mathcal{O}(m_{3/2})$, respectively. The lowest eigenvalue, which we identify with the pseudo goldstino mass m_ζ , can be calculated by keeping only the first two terms in Eq. (45),

$$m_\zeta = 2m_{3/2}\mathcal{Z}_1\mathcal{Z}_2, \quad (48)$$

where the suppression factor \mathcal{Z}_1 is given in Eq. (38) and

$$\mathcal{Z}_2^{-1} = 1 + \frac{\lambda^2 v^2 \sin^2 \beta}{\mu_2^2} + \frac{\lambda^2 v^2 \cos^2 \beta}{\mu_1^2}. \quad (49)$$

We see that the mass of the true pseudo goldstino is suppressed by a factor $\mathcal{Z}_1\mathcal{Z}_2$. We note that the lightness of the pseudo goldstino will be protected against quantum corrections by the $U(1)_R$ symmetry in the visible sector, as long as Dirac nature of the gauginos is kept.

The corresponding eigenvector can be found perturbatively in R -breaking effects. By writing $\vec{\Psi}_\zeta = \vec{\Psi}_\zeta^{(0)} + \vec{\Psi}_\zeta^{(1)} + \dots$, we find

$$\begin{aligned}\vec{\Psi}_\zeta^{(0)T} &= \sqrt{\mathcal{Z}_2} \left(\begin{array}{cc|cc|cc|c} 0 & 0 & 0 & 0 & 0 & 0 & \vec{\zeta}_{\tilde{X}}^{(0)T} & 1 \end{array} \right), \\ \vec{\Psi}_\zeta^{(1)T} &= \sqrt{\mathcal{Z}_2} \left(\begin{array}{c|cc|cc|c} \vec{\zeta}_\chi^{(1)T} & 0 & 0 & \vec{\zeta}_{\tilde{H}}^{(1)T} & 0 & 0 & 0 \end{array} \right).\end{aligned}\quad (50)$$

The two component vectors $\vec{\zeta}_{\tilde{X}}^{(0)}$, $\vec{\zeta}_\chi^{(1)}$ and $\vec{\zeta}_{\tilde{H}}^{(1)}$ are given respectively by

$$\vec{\zeta}_{\tilde{X}}^{(0)} = \begin{pmatrix} -\frac{\lambda v c_\beta}{\mu_1} \\ +\frac{\lambda v s_\beta}{\mu_2} \end{pmatrix}, \quad \vec{\zeta}_\chi^{(1)} = m_{3/2} \begin{pmatrix} -\frac{\sin \theta_W}{m_{\tilde{B}}} \\ \frac{\cos \theta_W}{m_{\tilde{W}}} \end{pmatrix} \mathcal{Z}_3, \quad \vec{\zeta}_{\tilde{H}}^{(1)} = m_\zeta \begin{pmatrix} \frac{\lambda v s_\beta}{\mu_2^2} \\ \frac{\lambda v c_\beta}{\mu_1^2} \end{pmatrix}, \quad (51)$$

where the dimensionless factor \mathcal{Z}_3 is defined by

$$\mathcal{Z}_3 = -m_Z s_\beta c_\beta \left(\frac{\lambda v}{\mu_2^2} - \frac{\lambda v}{\mu_1^2} \right) \frac{m_\zeta}{m_{3/2}} + \lambda v \left(\frac{s_\beta s_\gamma}{\mu_2} + \frac{c_\beta c_\gamma}{\mu_1} \right) \frac{m_V}{m_{3/2}}. \quad (52)$$

We see that at the leading order, only the xinos $\tilde{X}_{d,u}^0$ can mix with the would-be goldstino \tilde{X}_0 because of the $U(1)_R$ charge conservation; at the next order, the \tilde{X}_0 mixes with the fermions with $R = -1$, that is, the higgsinos $\tilde{H}_{d,u}^0$ and the gaugino Dirac partners $\chi_{B,W}$, since the gravitino mass $m_{3/2}$ has $R = -2$ in the sense of spurion.

In the sequestered limit, we have $\mathcal{Z}_1 \rightarrow 1$, $\mathcal{Z}_2 \rightarrow 1$ and $\mathcal{Z}_3 \rightarrow 0$, so that the eigenvalue and the eigenvector reduce to

$$m_\zeta \longrightarrow 2m_{3/2}, \quad \vec{\Psi}_\zeta^T \longrightarrow \left(\begin{array}{ccc|ccc|c} 0 & 0 & 0 & 0 & 0 & 0 & 1 \end{array} \right), \quad (53)$$

as is expected. In other cases, the mass eigenvalue m_ζ is much suppressed, and hence, the higgsino component $\vec{\zeta}_{\tilde{H}}^{(1)}$ is more suppressed than that of the gaugino Dirac partner, $\vec{\zeta}_\chi^{(1)}$:

$$\vec{\zeta}_{\tilde{H}}^{(1)} = \mathcal{O}\left(\frac{m_\zeta}{m_{\text{weak}}}\right), \quad \vec{\zeta}_\chi^{(1)} = \mathcal{O}\left(\frac{m_{3/2}}{m_{\text{weak}}}\right). \quad (54)$$

C. Implications on Higgs Phenomenology and Cosmology

Let us briefly discuss implications on the Higgs phenomenology and cosmology.

First, the invisible decay width of the 125 GeV Higgs into a pair of the pseudo goldstinos is highly suppressed due to the approximate R -symmetry. This is because fields with R -charge assignment other than +1 contain a small goldstino component with suppression factor $m_{2/3}/m_{\text{weak}}$. The Higgs decay into a pseudo goldstino and other neutralino is kinematically forbidden if the neutralino other than the pseudo goldstino are heavier than the lightest Higgs boson.

The second remark is concerned with the invisible cascade decay of the 125 GeV Higgs. If the scalar partner ϕ of the pseudo goldstino is so light that the decay mode of the Higgs into

a pair of ϕ is kinematically allowed, the cascade decay into gravitinos and pseudo goldstinos $h^0 \rightarrow \phi\phi \rightarrow \zeta\tilde{G}\zeta\tilde{G}$ dominates almost entire Higgs decay channel [6], which contradicts with recent LHC discovery [30]. In our case, the scalar goldstino ϕ receives a mass $\lambda^2 v^2$ from the visible sector SUSY breaking, in addition to the hidden-sector soft mass. Therefore, the invisible cascade decay of the Higgs is kinematically forbidden, if we assume $\lambda > 0.36$, or if the scalar pseudo goldstino has a soft mass larger than the half the Higgs mass.

Next we briefly discuss the cosmological constraints on the present model. Since we are supposing the low-scale SUSY breaking in the visible sector, the lightest observable-sector supersymmetric particle (LOSP) mainly decays into a pseudo goldstino rather than a gravitino. This fact helps us avoid the constraint from the BBN [11].

We should still worry about the overproduction problem, both for the gravitino and the pseudo goldstino. Let us first discuss the gravitino case. If gravitinos are thermally produced in the early Universe, the estimated abundance easily exceeds the limit $\Omega h^2 < 0.1$ for the gravitino heavier than $\mathcal{O}(100 \text{ eV})$. On the other hand, if gravitinos are not thermalized [31], the overproduction excludes the mass smaller than 10 keV, whereas the gravitino heavier than $\mathcal{O}(10 \text{ keV})$ can evade the overproduction if the reheating temperature is sufficiently low¹¹. Moreover, for the mass smaller than $\mathcal{O}(100 \text{ eV})$, the allowed region is much reduced to $m_{3/2} \lesssim 16 \text{ eV}$ due to a warm dark matter constraint [33, 34]. Putting it all together, we must assume that the gravitino mass satisfies either $m_{3/2} \lesssim 16 \text{ eV}$ or $10 \text{ keV} \lesssim m_{3/2}$, as was announced in Eq. (1)

In contrast, the pseudo goldstino couples a bit strongly to the MSSM particles and is expected to be thermalized. Recalling that its mass m_ζ is much suppressed compared to the gravitino mass $m_{3/2}$ as in Eq. (38), the pseudo goldstino will be lighter than 16 eV for the gravitino mass $m_{3/2} \gtrsim 10 \text{ keV}$; depending on the suppression factor \mathcal{Z}_1 , it may be possible that the pseudo goldstino evades the overclosure problem even when the gravitino mass is as large as $\mathcal{O}(10 \text{ MeV})$. If so, the pseudo goldstino can give only a small contribution to the relic density.

In the present model, the lightest supersymmetric particle (LSP) is the pseudo goldstino whereas the next-to-LSP (NLSP) is the gravitino, which can decay into a photon and a pseudo goldstino. In this situation, we should also be careful about the constraints from the gamma-ray line search from the Galactic Center region and the overproduction of the isotropic diffuse photon background; there is a stringent limit [35] on mono-energetic photons emitted from decaying particles. In our case, we expect that such decay of the gravitino is very suppressed since the pseudo goldstino mass eigenstate has very small gaugino and higgsino components. If this is really the case, the gravitino is sufficiently long-lived and the present model can evade the constraint from diffuse gamma-ray line search. A preliminary consideration shows that only the gravitino heavier than $\mathcal{O}(100 \text{ MeV})$ would be excluded even if the mixing between the photino and the pseudo goldstino is of order of $\mathcal{O}(m_{3/2}/m_{\text{weak}})$.

¹¹ If the reheating temperature is low, the thermal leptogenesis is difficult to be achieved. A possibility of electroweak baryogenesis in R -symmetric models as was recently discussed in Ref. [32].

V. CONCLUSION AND DISCUSSION

We have studied the $U_R(1)$ symmetric extension of SUSY standard model which contains Dirac gauginos and extended Higgs sector, assuming that the Dirac gaugino masses are induced from hidden sector SUSY breaking while the extended Higgs sector incorporates the visible SUSY breaking. The $U(1)_R$ symmetry in the visible sector is broken solely by the minimal coupling to supergravity, so that $U(1)_R$ breaking effects are proportional to the gravitino mass. After the EWSB, the R -charged Higgs fields develop nonzero VEVs due to the tadpole terms induced from the $U(1)_R$ breaking. The VEVs are proportional to the gravitino mass and hence can be very small when the gravitino mass is much smaller than the weak scale.

We have then shown the generation of small neutrino masses through the $U(1)_R$ breaking effects. With our flipped $U(1)_R$ charge assignment, the right-handed neutrinos couple to the R -partner Higgs field X_u that develops the very small VEV. Therefore the small neutrino masses can naturally be obtained even if the neutrino Yukawa couplings are of order unity, provided that the gravitino mass is as low as 1–10 eV. We have also examined how our generation mechanism of neutrino mass is constrained from the proton decay induced by the dimension five operator, $U^c D^c D^c N^c$, allowed by our $U(1)_R$ charge assignment. Interestingly, the amplitude for the dominant decay mode involves the neutrino Yukawa coupling, instead of the top Yukawa. Therefore the constraint from proton stability is severer for a larger neutrino Yukawa coupling, which implies a small gravitino mass for a fixed neutrino mass. We have estimated the lifetime of protons and found that the gravitino should be heavier than 1 keV if we require the coefficient C_5 of the dimension five operator to be of order unity and the cutoff Λ_{cutoff} to be the Planck scale. Actually, the gravitino mass between $10 \text{ eV} < m_{3/2} < 10 \text{ keV}$ is cosmologically excluded [33, 34, 36], and hence our model for the neutrino mass generation evade the constraint from proton decay and is cosmologically safe if gravitino is as heavy as 10–100 keV. We have also suggested how the constraint from the proton stability can be relaxed if we adopt the $U(1)$ flavor symmetry for generating the hierarchical structure of Yukawa matrices.

Another significant effect of the $U(1)_R$ breaking is the mass and the mixing of the pseudo goldstino, which is Nambu-Goldstone fermion of the visible SUSY breaking and is massless in the $U(1)_R$ symmetric limit. We have analyzed the neutralino mass matrix that incorporates the Dirac mass parameters of gauginos and higgsinos, and obtained to the lowest order in R -breaking the smallest mass eigenvalue and the corresponding mixing angles to other neutralinos. According to the general argument, the mass of the pseudo goldstino in the sequestered limit is twice the gravitino mass. In our case, the pseudo goldstino has a coupling to the Higgs fields, and consequently, its mass is proportional to but suppressed from the gravitino mass. In our concrete calculation, in which the pseudo goldstino gets a mass from the contact term in the Kähler potential, the suppression factor is very small if the cutoff scale Λ_0 of the contact term is much larger than the weak scale. For the mixings, we have found that the pseudo goldstino have suppressed contaminations of gauginos and higgsinos.

As we mentioned above, the gravitino as heavy as 100 keV is cosmologically safe and

also satisfies the proton stability. In this case, the gravitino is the NLSP while the LSP is the pseudo goldstino, which can be as light as 10 eV, or even much lighter. It is then tempting to speculate that the gravitino could be a part of cold dark matter, and that the pseudo goldstino could play a role of dark radiation. To examine such possibilities, we should investigate the properties of the pseudo goldstino more extensively, for instance, by including the anomaly mediated gaugino masses and other quantum corrections to the mass matrix. This will be important also for a quantitative study of diffuse gamma-ray line.

Although we have focused in the present work on the constraints from the proton stability and also on some cosmological ones, there remain many issues to be discussed: the constraints from EW precision measurements and an implementation of the 125 GeV Higgs should be examined more intensively: phenomenological impacts of large neutrino Yukawa couplings as in other “neutrinophilic” models may deserve further study in our $U(1)_R$ symmetric model. The assumptions behind our present setup should also be elucidated: a partial list includes the origin of the hidden sector SUSY breaking, the UV completion of the visible SUSY breaking sector and the embedding into a semi-simple gauge group. In particular, the smallness of the gravitino mass should be examined in a consistent framework, although we have treated it simply as a free parameter. In this respect, it may be possible that the smallness of gravitino mass as well as the absence of certain operators can be explained in the (warped) extra-dimensional setup along the line of Ref. [37].

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Appendix A: Baryon and Lepton Number Violating Operators

In this appendix, we systematically study baryon and/or lepton number violating operators allowed by the gauge and $U(1)_R$ symmetries of the present model.

We first categorize the fields with the same gauge quantum numbers as follows:

$$S \in \{X_0, N^c\} , \quad \Phi_u \in \{H_u, X_u\} , \quad \Phi_d \in \{H_d, X_d, L\} . \quad (\text{A1})$$

The operators allowed by the gauge symmetries can be classified according to their mass dimensions as

$$W_2 = S , \quad (\text{A2})$$

$$W_3 = SW_2 + \Phi_u \Phi_d , \quad (A3)$$

$$W_4 = SW_3 + (Q\Phi_d)D^c + (\Phi_d\Phi_d)E^c + U^c D^c D^c , \quad (A4)$$

$$W_5 = SW_4 + (QQ)(Q\Phi_d) + U^c U^c D^c E^c + (Q\Phi_d)U^c E^c \\ + (\Phi_u\Phi_d)(\Phi_u\Phi_d) + (QQ)U^c D^c , \quad (A5)$$

where we suppress coefficients and flavor indices for simplicity. We also omit operators involving the adjoint chiral fields since they have no effect on proton decay.

It is straightforward to see that all of the operators violating either the baryon or lepton number in W_2 , W_3 and W_4 are forbidden by the R -symmetry: this includes the usual trilinear terms, $U^c D^c D^c$, $QD^c L$, $LE^c L$, and the bilinear terms $H_u L$. We also note that the mass terms $X_0 N^c$ as well as $N^c N^c$ are forbidden by the $U(1)_R$. Therefore we can define the conserved baryon and lepton numbers at the renormalizable level.

There are three types of dimension five operators that violate both the baryon and lepton numbers: $(QQ)(QL)$, $U^c U^c D^c E^c$ and $N^c U^c D^c D^c$. Among them, the first two are forbidden by the $U(1)_R$, whereas the last one is $U(1)_R$ symmetric in our flipped assignment. As a result, we are left with a unique operator $N^c U^c D^c D^c$, which is the baryon and lepton number violating operator in the present model.

As we mentioned in Sec. III A, $(LH_u)(LH_u)$ operators are allowed, but they give a negligible contribution to neutrino masses if the cutoff Λ_{cutoff} is sufficiently large. We also note that there are dimension five operators, $X_0^2 N^c N^c$ and $(X_u X_d)N^c N^c$, which are R -invariant and lepton number violating. After the R -Higgses get the VEVs, these operators lead to Majorana masses for right-handed neutrinos, of order $m_{3/2}^2/\Lambda$, which can safely be neglected (unless $m_{3/2} \sim 100$ TeV).

Appendix B: $U(1)$ Flavor Symmetry and Proton Stability

As we see in sec.III, the model is severely constrained from proton decay if the gravitino is very light, $m_{3/2} < 16$ eV, especially for large $\tan\beta$. In this appendix, we consider a model with Froggatt-Nielsen $U(1)$ flavor symmetry to relax the constraint.

Let Θ be a Froggatt-Nielsen (FN) field with $U(1)_F$ charge -1 , which is assumed to develop a nonzero VEV $\langle\Theta\rangle$. We also denote the $U(1)_F$ charge of the fields Φ_i by the lower case letter ϕ_i . The Yukawa terms in Eq. (18) should be multiplied by a suitable power of the FN field Θ ; for instance, the neutrino Yukawa term is generated from the $U(1)_F$ -invariant operator

$$W = f_N^{ij} \left(\frac{\Theta}{\Lambda_{\text{cutoff}}} \right)^{n_i^c + \ell_j + X_u} N_i^c L_j X_u ,$$

where f_N^{ij} is a coefficient of order 1 and Λ_{cutoff} is the cutoff. After we substitute the FN field by its VEV, we have $y_N^{ij} = f_N^{ij} \varepsilon^{\nu_i^c + \ell_j + X_u}$. We will identify $\varepsilon \equiv \langle\Theta\rangle/\Lambda_{\text{cutoff}}$ with the Wolfenstein parameter $\varepsilon = 0.22$. Similarly, we have

$$y_U^{ij} = f_U^{ij} \varepsilon^{q_i + u_j^c + h_u} , \quad y_D^{ij} = f_D^{ij} \varepsilon^{q_i + d_j^c + h_d} , \quad y_E^{ij} = f_E^{ij} \varepsilon^{e_i^c + l_j + h_d} . \quad (B1)$$

The $U(1)_F$ charge assignment can be restricted by the following requirements. First, from the intergenerational mixings of quarks, $V_{us} \sim \varepsilon$, $V_{cb} \sim \varepsilon^2$ and $V_{ub} \sim \varepsilon^3$, we obtain the following relations

$$q_1 = q_3 + 3, \quad q_2 = q_3 + 2. \quad (\text{B2})$$

Next, the quark mass hierarchy can be parameterized as $m_u : m_c : m_t \sim \varepsilon^8 : \varepsilon^4 : 1$ and $m_d : m_s : m_b \sim \varepsilon^4 : \varepsilon^2 : 1$, while the mass hierarchy of charged leptons is $m_e : m_\mu : m_\tau \sim \varepsilon^5 : \varepsilon^2 : 1$. These mass hierarchies can be reproduced, for instance, if we take

$$\begin{aligned} u_1^c &= u_3^c + 5, & d_1^c &= d_3^c + 1, & \ell_1 + e_1^c &= \ell_3 + e_3^c + 5, \\ u_2^c &= u_3^c + 2, & d_2^c &= d_3^c, & \ell_2 + e_2^c &= \ell_3 + e_3^c + 2. \end{aligned} \quad (\text{B3})$$

The FN charges of the third generation of quarks and leptons can be constrained, if we require the top Yukawa coupling of order unity, $y_t \sim 1$, and also the bottom-tau unification, $m_b \sim m_\tau$, which gives

$$q_3 + u_3^c + h_u = 0, \quad q_3 + d_3^c = e_3^c + \ell_3, \quad (\text{B4})$$

In addition, we require for definiteness that the supersymmetric Higgs mass terms have null charge,

$$x_d + h_u = 0, \quad x_u + h_d = 0. \quad (\text{B5})$$

Note that the FN charges of H_u and H_d can be chosen independently since there is no $\mu H_u H_d$ term in the present $U(1)_R$ symmetric model.

Now, we discuss the suppression of the dimension five operator (22) by the FN mechanism. The coefficient $C_5 = C_5^{1123}$ of the operator $U_1^c D_1^c D_2^c N_3^c$ receives a suppression of $C_5 \sim \varepsilon^{u_1^c + d_1^c + d_2^c + n_3^c}$. The FN charge n_3^c of the right-handed neutrino can be eliminated if we use the relation (26)

$$y_\nu^{ij} = f_\nu^{ij} \varepsilon^{n_i^c + l_j + x_u} = \frac{m_\nu^{ij}}{m_{3/2}} \left(\frac{\mu_2^2 + m_{X_u}^2}{\mu_2 v \cos \beta} \right). \quad (\text{B6})$$

Combining Eqs. (B2)–(B5) with Eq. (B6), we obtain

$$C_5 = C_5^{(0)} \varepsilon^{u_1^c + d_1^c + d_2^c + n_3^c} = C_5^{(0)} \varepsilon^{6+Q} \left(\frac{m_\nu}{m_{3/2}} \right) \left(\frac{\mu_2^2 + m_{X_u}^2}{\mu_2 v \cos \beta} \right), \quad (\text{B7})$$

where

$$Q \equiv h_d - h_u + d_3^c + e_3^c - 2q_3. \quad (\text{B8})$$

We see that the coefficient C_5 is suppressed for several reasons: the suppression factor ε^6 comes from the FN charges of the 1st and 2nd generations of quarks involved in the proton decay. Another factor ε^Q expresses the dependence on the FN charges of the Higgses and the 3rd generation of quarks and leptons. For instance, if we assign a negative FN charge to the MSSM Higgs H_u , the up-type quarks have positive FN charges, which suppress the dimension

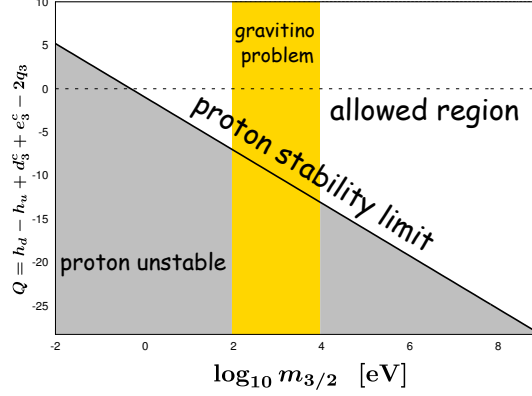


FIG. 4. The constraint on the FN charge Q . The parameters are the same as in Fig. 3.

five operator. Notice that this can be done without suppressing the supersymmetric mass terms since we can assign the opposite FN charge to the R -partner X_d .

It should be noticed that there is yet another suppression in Eq. (B7): the coefficient C_5 is more suppressed for larger gravitino mass. This dependence can be understood as follows. Assigning a positive FN charge $n^c > 0$ to the right-handed neutrinos suppresses not only the coefficient C_5 of the dimension five operator but also the neutrino Yukawa coupling y_ν ; $y_\nu \sim \varepsilon^{n^c}$ and $C_5 \sim \varepsilon^{n^c}$, so that the proton lifetime scales as $\tau_p \sim \varepsilon^{-4n^c}$. For a fixed value of the neutrino mass, the suppression of the neutrino Yukawa coupling is translated into the dependence on $m_{3/2}$ in Eq. (B7).

Putting it all together, a similar analysis as in Sec. III B gives an estimate of the proton lifetime as

$$\tau_p = \frac{3.7 \times 10^{34} \text{ yr}}{\varepsilon^{2Q}} \left(\frac{\Lambda_{\text{cutoff}}}{2.4 \times 10^{18} \text{ GeV}} \right)^2 \left(\frac{m_{3/2}}{1 \text{ eV}} \right)^4 \left(\frac{\cos^2 \beta}{0.1} \right)^3, \quad (\text{B9})$$

where we have set the coefficient to be $C_5^{(0)} = 1$. Fig. 4 shows how the constraint from the proton decay can be satisfied for a moderate choice of the FN charge Q . We see that the proton decay constraint can be satisfied even for $m_{3/2} = 1 \text{ eV}$ if we take $Q = 0$ and $\tan \beta = 3$. In this way, the present model combined with the $U(1)_F$ flavor symmetry evades the proton decay constraint in a wider region of parameter space.

The situation is quite different from that in the MSSM case [29], in which a negative FN charge of the Higgses would imply unacceptably small μ term. In fact, the analysis in Ref. [29] shows that the FN charge of the μ term, *i.e.*, the sum of the FN charges of H_u and H_d , are tightly constrained since both of LLLL and RRRR operators (31) should be suppressed simultaneously. In our case, we have only RRRR type operator since LLLL operator is forbidden by the $U(1)_R$ symmetry.

[1] G. D. Kribs, E. Poppitz, and N. Weiner, Phys.Rev. **D78**, 055010 (2008), 0712.2039.

- [2] P. J. Fox, A. E. Nelson, and N. Weiner, JHEP **0208**, 035 (2002), hep-ph/0206096.
- [3] G. D. Kribs and A. Martin, Phys.Rev. **D85**, 115014 (2012), 1203.4821.
- [4] K. Izawa, Y. Nakai, and T. Shimomura, JHEP **1103**, 007 (2011), 1101.4633.
- [5] K. Benakli, M. D. Goodsell, and F. Staub (2012), 1211.0552.
- [6] D. Bertolini, K. Rehermann, and J. Thaler, JHEP **1204**, 130 (2012), 1111.0628.
- [7] N. Arkani-Hamed, L. J. Hall, H. Murayama, D. Tucker-Smith, and N. Weiner, Phys.Rev. **D64**, 115011 (2001), hep-ph/0006312.
- [8] J. March-Russell and S. M. West, Phys.Lett. **B593**, 181 (2004), hep-ph/0403067.
- [9] K. Rehermann and C. M. Wells (2011), 1111.0008.
- [10] R. Davies and M. McCullough, Phys.Rev. **D86**, 025014 (2012), 1111.2361.
- [11] C. Cheung, Y. Nomura, and J. Thaler, JHEP **1003**, 073 (2010), 1002.1967.
- [12] R. Argurio, Z. Komargodski, and A. Mariotti, Phys.Rev.Lett. **107**, 061601 (2011), 1102.2386.
- [13] S. Shirai, M. Yamazaki, and K. Yonekura, JHEP **1006**, 056 (2010), 1003.3155.
- [14] S. Ray (2007), 0708.2200.
- [15] C. Cheung, F. D’Eramo, and J. Thaler, Phys.Rev. **D84**, 085012 (2011), 1104.2598.
- [16] E. Ma, Phys.Rev.Lett. **86**, 2502 (2001), hep-ph/0011121.
- [17] E. Ma and M. Raidal, Phys.Rev.Lett. **87**, 011802 (2001), hep-ph/0102255.
- [18] S. M. Davidson and H. E. Logan, Phys. Rev. D **80**, 095008 (2009), URL <http://link.aps.org/doi/10.1103/PhysRevD.80.095008>.
- [19] T. Morozumi, H. Takata, and K. Tamai, Phys. Rev. D **85**, 055002 (2012), URL <http://link.aps.org/doi/10.1103/PhysRevD.85.055002>.
- [20] N. Haba and T. Horita, Physics Letters B **705**, 98 (2011), ISSN 0370-2693, URL <http://www.sciencedirect.com/science/article/pii/S0370269311012032>.
- [21] N. Haba and O. Seto, Prog.Theor.Phys. **125**, 1155 (2011), 1102.2889.
- [22] K.-Y. Choi and O. Seto, Phys. Rev. D **86**, 043515 (2012), URL <http://link.aps.org/doi/10.1103/PhysRevD.86.043515>.
- [23] F. Wang, W. Wang, and J. M. Yang, Europhys.Lett. **76**, 388 (2006), hep-ph/0601018.
- [24] M. Sher and C. Triola, Phys. Rev. D **83**, 117702 (2011), URL <http://link.aps.org/doi/10.1103/PhysRevD.83.117702>.
- [25] S. Zhou, Phys.Rev. **D84**, 038701 (2011), 1106.3880.
- [26] T. Goto and T. Nihei, Phys.Rev. **D59**, 115009 (1999), hep-ph/9808255.
- [27] S. Aoki et al. (JLQCD Collaboration), Phys.Rev. **D62**, 014506 (2000), hep-lat/9911026.
- [28] K. Kobayashi, M. Earl, Y. Ashie, J. Hosaka, K. Ishihara, Y. Itow, J. Kameda, Y. Koshio, A. Minamino, C. Mitsuda, et al. (Super-Kamiokande Collaboration), Phys. Rev. D **72**, 052007 (2005), URL <http://link.aps.org/doi/10.1103/PhysRevD.72.052007>.
- [29] M. Kakizaki and M. Yamaguchi, JHEP **0206**, 032 (2002), hep-ph/0203192.
- [30] P. P. Giardino, K. Kannike, M. Raidal, and A. Strumia (2012), 1207.1347.
- [31] M. Endo, F. Takahashi, and T. Yanagida, Phys.Rev. **D76**, 083509 (2007), 0706.0986.
- [32] R. Fok, G. D. Kribs, A. Martin, and Y. Tsai (2012), 1208.2784.
- [33] M. Viel, J. Lesgourgues, M. G. Haehnelt, S. Matarrese, and A. Riotto, Phys.Rev. **D71**, 063534 (2005), astro-ph/0501562.
- [34] A. Boyarsky, J. Lesgourgues, O. Ruchayskiy, and M. Viel, JCAP **0905**, 012 (2009), 0812.0010.

- [35] H. Yuksel and M. D. Kistler, Phys.Rev. **D78**, 023502 (2008), 0711.2906.
- [36] K. Ichikawa, M. Kawasaki, K. Nakayama, T. Sekiguchi, and T. Takahashi, JCAP **0908**, 013 (2009), 0905.2237.
- [37] Z. Chacko, P. J. Fox, and H. Murayama, Nucl.Phys. **B706**, 53 (2005), hep-ph/0406142.